

# A possible connection between neutrino mass generation and the lightness of a NMSSM pseudoscalar

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## Abstract

One of the interesting properties of the NMSSM is that it can accommodate a light pseudoscalar of order 10 GeV. However, such scenarios are challenged by several experimental constraints, especially those related to the fermionic decays of the pseudoscalar. In this Letter, we extend the NMSSM field content by two gauge singlets, with lepton numbers +1 and −1. This serves the twin purpose of generating neutrino masses via the inverse seesaw mechanism and keeping the option of a very light pseudoscalar experimentally viable by opening dominant invisible decay channels of the pseudoscalar which help it evade the existing bounds.

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We consider the Next to Minimal Supersymmetric Standard Model (NMSSM) which provides a natural solution to the so-called  $\mu$  problem through the introduction of a new gauge singlet superfield  $\hat{S}$  in the superpotential. In the NMSSM, the  $\mu$  parameter is linked to the vacuum expectation value (VEV) of the scalar component of  $\hat{S}$  whose size is of the order of the supersymmetry breaking scale [1]. Since it permits a scale invariant superpotential, the NMSSM is the simplest supersymmetric generalization of the Standard Model (SM) in which the supersymmetry breaking scale is the only mass scale in the Lagrangian. Moreover, by providing an additional tree level contribution to the quartic term of the scalar potential it can ameliorate the ‘little hierarchy problem’ of the MSSM, which is related to the requirement of a large ( $\gg M_Z$ ) soft supersymmetry breaking mass that can push the mass of the lightest neutral Higgs beyond the LEP-2 limit of 114 GeV [2]. Furthermore, the NMSSM can admit a very light CP-odd Higgs boson ( $m_{A_1} \sim 1 - 10$  GeV) [3, 4]. The main objective of this Letter is to improve the experimental viability of such a light pseudoscalar by providing it with a dominant invisible decay mode in a minimal extension of the NMSSM which contains a source of lepton number violation that also yields an acceptable neutrino mass.

In the NMSSM, the lightest CP-odd physical scalar  $A_1$  can be decomposed as

$$A_1 \equiv \cos \theta_A A_{\text{MSSM}} + \sin \theta_A A_S, \quad (1)$$

where  $A_{\text{MSSM}}$  is the MSSM part of the CP-odd scalar, which arises solely from the NMSSM Higgs doublets, and  $A_S$  is the part that arises from the new singlet superfield  $\hat{S}$ . It is the singlet admixture,

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i.e. the  $\sin\theta_A$  projection, that allows the NMSSM pseudoscalar to be much lighter than what it could have been in the MSSM. On the other hand, if  $A_1$  is very light then its detection crucially depends on its couplings to quarks and leptons, which depend on  $\cos\theta_A$ . These couplings can be extracted from the following part of the Lagrangian [5]:

$$\mathcal{L}_{A_1 f \bar{f}} = X_{u(d)} \frac{gm_f}{2M_W} \bar{f} \gamma_5 f A_1, \quad (2)$$

where  $g$  is the SU(2) gauge coupling,  $X_d(X_u) = \cos\theta_A \tan\beta$  ( $\cos\theta_A \cot\beta$ ) for down-type (up-type) fermions,  $\tan\beta \equiv v_u/v_d$  with  $v_u$  and  $v_d$  denoting the up- and down-type Higgs VEVs.

A light pseudoscalar is phenomenologically interesting mainly for two reasons:

(i) The direct search limit of 114 GeV on the mass of the SM-like Higgs (a slightly smaller limit of 93 GeV for the MSSM neutral Higgs) is obtained from its non-observation at the highest energy run at LEP-2, where the Higgs was expected to be produced from gauge interactions (full strength  $ZZh$  coupling) and decay dominantly ( $\sim 75\%$ ) into  $b\bar{b}$  final states. In the NMSSM, there could be two important changes:

- The lightest CP-even Higgs ( $h$ ) may have a large singlet component, which leads to a significant dilution of the  $ZZh$  coupling. The Higgs production cross section will then be reduced, and hence the lower limit on  $m_h$  will be relaxed.
- $h$  may dominantly decay into a pair of  $A_1$ , with each  $A_1$  decaying into  $f\bar{f}$ , where  $f$  is  $b$ ,  $\tau$ ,  $c$ , or  $\mu$ , depending on the kinematic thresholds. Therefore, the existing LEP-2 Higgs search strategy from  $2b$  final states would fail as one should look for  $4f$  final states. In fact, a reanalysis of the LEP-2 data by the LEP Collaboration has already put constraints on  $\frac{\sigma(e^+e^- \rightarrow Zh)}{\sigma_{\text{SM}}(e^+e^- \rightarrow Zh)} \times \text{Br}(h \rightarrow A_1 A_1) \times \text{Br}(A_1 \rightarrow f\bar{f})^2$ , where  $f = b$  or  $\tau$  depending on kinematics thresholds [6, 7]. Similarly, for  $m_{A_1} < 2m_\tau$ , upper limits have been placed on  $\sigma(p\bar{p} \rightarrow hX) \times \text{Br}(h \rightarrow A_1 A_1) \times \text{Br}(A_1 \rightarrow \mu^+ \mu^-)^2$  by the D0 collaboration at Fermilab Tevatron [8]. The upshot is that the lower limit on  $m_h$  is slightly weakened, and a smaller value of  $m_h$  (e.g. 105 GeV) is in fact preferred by electroweak precision tests [9].

(ii) If the lightest supersymmetric particle (LSP) happens to be very light (a few GeV), then a light  $A_1$  offers the possibility of  $s$ -channel LSP pair-annihilation into an on-shell  $A_1$ . This resonance channel has a special significance when one attempts to account for the observed dark matter relic abundance. It has recently been shown that a light LSP of mass  $\sim 10$  GeV can have interesting consequences in the context of the recent DAMA/CoGeNT results [10].

Let us now briefly discuss the existing bounds on the mass of the light pseudoscalar. The constraints on  $X_d$ , defined in Eq. (2), for  $m_{A_1}$  approximately in the range of 1 to 10 GeV have been summarized in [4, 5]. Measurements of  $\Delta M_{d,s}$ ,  $\text{Br}(\bar{B} \rightarrow X_s \gamma)$ ,  $\text{Br}(B^+ \rightarrow \tau^+ \nu_\tau)$ , and particularly,  $\text{Br}(\bar{B}_s \rightarrow \mu^+ \mu^-)$  severely constrain  $m_{A_1}$  [11]. The rates of these processes primarily depend on the choice of  $\tan\beta$  and the soft supersymmetry breaking trilinear term  $A_t$ , and the constraints are in general weaker when these parameters are small. Values of  $m_{A_1}$  between 1 GeV and  $m_b$  are generally disfavored from  $B$ -meson data [4]. Constraints on  $m_{A_1}$  also arise from radiative  $\Upsilon$  decays [12, 13, 14], namely,  $\Upsilon(nS) \rightarrow \gamma A_1$ , with  $A_1 \rightarrow \mu^+ \mu^- (\tau^+ \tau^-)$  (further investigated and reviewed in [5]). Severe constraints also arise as a consequence of  $\eta_b - A_1$  mixing [15, 16, 17]. The different  $m_{A_1}$  windows which are sensitive to different processes are listed in Table 1. The table also shows the ranges where the LEP (ALEPH [7] and OPAL [18]) constraints are applicable. The origin of all these constraints can be traced to the visible decay modes of  $A_1$ .

Processes	$m_{A_1} < 2m_\tau$	$[2m_\tau, 9.2 \text{ GeV}]$	$[9.2 \text{ GeV}, M_{\Upsilon(1S)}]$	$[M_{\Upsilon(1S)}, 2m_B]$
$\Upsilon \rightarrow \gamma A_1 \rightarrow \gamma + (\mu^+ \mu^-, gg, s\bar{s})$	✓	×	×	×
$\Upsilon \rightarrow \gamma A_1 \rightarrow \gamma \tau^+ \tau^-$	×	✓	×	×
$A_1$ - $\eta_b$ mixing	×	×	✓	×
$e^+ e^- \rightarrow Z + 4\tau$ (ALEPH)	×	✓	✓	✓
$e^+ e^- \rightarrow b\bar{b} \tau^+ \tau^-$ (OPAL)	×	×	✓	✓

Table 1: *Different processes constraining different  $m_{A_1}$  windows. The “✓” symbol in a given entry attests the existence of important or meaningful constraints from a given process, while the “×” symbol implies otherwise.*

However, the situation may dramatically change if  $A_1$  has dominant invisible decay modes. Its decay into a pair of stable neutralinos (if kinematically possible) is one such example. The BABAR Collaboration [19] at the PEP-II  $B$ -factory has, however, searched for radiative  $\Upsilon$ -decays where a large missing mass is accompanied by a monochromatic photon, and from its non-observation has set a (preliminary) 90% C.L. upper limit on  $\text{Br}(\Upsilon(3S) \rightarrow \gamma A_1) \times \text{Br}(A_1 \rightarrow \text{invisible})$  at  $(0.7 - 31) \times 10^{-6}$  for  $m_{A_1}$  in the range of 3 to 7.8 GeV.

In this Letter we further explore the possibility of invisible decay channels that would allow a light  $A_1$  escape detection even outside the range of 3 to 7.8 GeV. We show that if we extend the NMSSM by two additional gauge singlets with non-vanishing lepton numbers, they would not only provide a substantial invisible decay channel of  $A_1$  but, as a bonus, would also generate small neutrino masses through lepton number violating ( $\Delta L = 2$ ) interactions. The visible decay branching ratios of  $A_1$  would then be reduced. As a result, the constraints on  $X_d$  would be weakened. A light  $A_1$  can then be comfortably accommodated.

In the framework of the NMSSM, neutrino masses and mixings can be generated via different mechanisms, either under the assumption of  $R$ -parity conservation or  $R$ -parity violation (RpV). In the latter case, light neutrino masses/mixing data can be successfully accommodated via the inclusion of explicit trilinear and/or bilinear RpV terms [20], or through spontaneous violation of lepton number in the presence of right-handed neutrino superfields (in addition to the NMSSM singlet  $\hat{S}$ ) [21]. Alternatively, one can employ a standard seesaw mechanism in a  $R$ -parity conserving setup by adding three gauge singlet neutrino superfields  $\hat{N}_i$  to the NMSSM particle content [22]. In this case, the light neutrino masses originate from the tiny Yukawa couplings and (dynamically generated) TeV scale Majorana mass terms.

Here, we follow none of the above paths. We rather implement the “inverse seesaw” mechanism [23] by adding two gauge singlet superfields with opposite lepton numbers (+1 and −1). During this implementation we assume that  $R$ -parity is conserved, i.e. we do not admit any  $\Delta L = 1$  term in the superpotential. To appreciate the advantages of the inverse seesaw over the standard one we look at the difference, from the point of view of effective operators, between the properties of the dimension-5 (d-5) Weinberg operator which is lepton number *violating* and the dimension-6 (d-6) operator which is lepton number *conserving*. In the d-5 case, the light neutrino mass is given by the seesaw formula  $m_\nu \sim m_D^2/M$ , where  $m_D = fv$  is a Dirac mass, with  $v$  as the electroweak VEV and  $f$  as a generic Yukawa coupling. The source of lepton number violation is the Majorana mass  $M$ . If we demand the Yukawa coupling  $f$  to be order one, then for  $m_\nu$  to be  $\sim 1$  eV,  $M$  has to be close to the gauge coupling unification scale. On the other hand, the d-6 operator goes as  $1/M^2$ , but since

this operator conserves lepton number, its coefficient is not related to that of the d-5 operator [24]. Therefore, for the d-6 case, one needs a separate source of lepton number violation to generate light neutrino mass. The generic form of this mass is  $m_\nu \sim (m_D^2/M^2)\mu$ , where  $\mu$  is a lepton number violating ( $\Delta L = 2$ ) mass parameter. In this case, one can comfortably keep the fundamental scale  $M$  close to TeV, i.e. within the LHC reach, and yet choose  $f$  to be order unity which can trigger large lepton flavor violation. Here, the lightness of the neutrino mass (eV, or even lighter) is due to the smallness of  $\mu$ . Having this small dimensionful term in the Lagrangian is technically natural in the sense of 't Hooft [25], as in the limit of vanishing  $\mu$  one recovers the lepton number symmetry. The structure of light neutrino mass obtained via inverse seesaw conforms to the principle of ‘low fundamental scale, large Yukawa coupling and light neutrino mass’ [26].

In our model, the superpotential is given by

$$W = W_{\text{NMSSM}} + W', \quad \text{where,} \quad (3)$$

$$W_{\text{NMSSM}} = f_{ij}^d \hat{H}_d \hat{Q}_i \hat{D}_j + f_{ij}^u \hat{H}_u \hat{Q}_i \hat{U}_j + f_{ij}^e \hat{H}_d \hat{L}_i \hat{E}_j + \lambda_H \hat{S} \hat{H}_d \hat{H}_u + \frac{\kappa}{3} \hat{S}^3, \quad (4)$$

$$W' = f_{ij}^\nu \hat{H}_u \hat{L}_i \hat{N}_j + (\lambda_N)_i \hat{S} \hat{N}_i \hat{X}_i + \mu_{X_i} \hat{X}_i \hat{X}_i. \quad (5)$$

In the above expressions,  $\hat{H}_d$  and  $\hat{H}_u$  are the down- and up-type Higgs superfields;  $\hat{Q}_i$  and  $\hat{L}_i$  denote the SU(2) doublet quark and lepton superfields;  $\hat{U}_i$  ( $\hat{D}_i$ ) and  $\hat{E}_i$  are the SU(2) singlet up (down)-type quark superfields and the charged lepton superfields, respectively. We have denoted the Yukawa couplings by  $f$  with appropriate flavor ( $u, d, e, \nu$ ) and generation indices ( $i, j = 1, 2, 3$ ).  $\hat{S}$  is the singlet superfield already present in the minimal NMSSM. Besides, we have added two more gauge singlets  $\hat{N}$  and  $\hat{X}$ , for each generation, which carry lepton numbers  $L = -1$  and  $L = +1$ , respectively. In our formulation, even though  $L$  is not a good quantum number because of the presence of a non-vanishing  $\mu_X$ ,  $(-1)^L$  is still a good symmetry. We have written the  $\hat{N} \hat{X}$  and  $\hat{X} \hat{X}$  terms in a generation diagonal basis without any loss of generality. Once the scalar component of  $\hat{S}$  acquires a VEV ( $v_S$ ), not only the conventional  $\mu$ -term is generated with  $\mu = \lambda_H v_S$ , a lepton number *conserving* mass term  $M_N \Psi_N \Psi_X$  is generated as well, with  $M_N = \lambda_N v_S$ . One more lepton number *conserving* mass term  $m_D \Psi_\nu \Psi_N$  emerges with  $m_D = f^\nu v_u$ .

The crucial term relevant for inverse seesaw is the  $\Delta L = 2$  term involving  $\mu_X$ , which is the only mass dimensional term in the superpotential. We assume that the  $Z_3$  symmetry of the superpotential is absent only in this term. We treat  $\mu_X$  as an extremely tiny *effective* mass parameter generated by some unknown dynamics. Its smallness would eventually decide the lightness of the light neutrino. We make a few observations at this stage:

(i) The superpotential treats the two singlets  $\hat{N}$  and  $\hat{X}$  differently in the sense that it yields a  $\mu_X \Psi_X \Psi_X$  Majorana mass term ( $\Delta L = 2$ ) but does not lead to a similar  $\mu_N \Psi_N \Psi_N$  term. This discrimination requires further qualification. A generic superpotential with  $(-1)^L$  parity should have included the latter term. Both  $\mu_N$  and  $\mu_X$  can be naturally small, as their absence enhances the symmetry of the Lagrangian. But the important thing to note is that the magnitude of  $\mu_X$  (and not that of  $\mu_N$ ) controls the size of the light neutrino mass [27, 28]. In view of this, for the sake of simplicity, we have assumed  $\mu_N = 0$ .

(ii) Two questions naturally arise here. First, although we have put  $\mu_X$  by hand and claimed that it is tiny, is it possible to dynamically generate a small  $\mu_X$  starting from a superpotential which does not *a priori* contain any mass dimensional term? Second, is it possible to realize  $\mu_N \ll \mu_X$  in a sensible model?

To address the first question, we admit that with the particle content of our model it is not possible to provide a natural solution for a small  $\mu_X$ . One possibility could have been to start with a trilinear  $\lambda \hat{S} \hat{X} \hat{X}$  term in the superpotential of Eq. (5), which would lead to  $\mu_X = \lambda v_S$ . Since  $v_S \sim v$ , the requirement to produce the correct light neutrino mass would then compel us to take  $\lambda \sim 10^{-11}$ . Such a small dimensionless trilinear coupling would be against the spirit of inverse seesaw mechanism as illustrated earlier. However, by extending the particle content of the model it is possible to dynamically generate a small  $\mu_X$ . In this context, we recall that in the original inverse seesaw formulation [23], the smallness of  $\mu_X$  was attributed to the supersymmetry breaking effects in a superstring inspired  $E_6$  scenario. It is also possible to keep  $\mu_X$  small by relating it to a tiny VEV generated dynamically. An analysis along this line was carried out in an extended version of the NMSSM [29], where the origin of small VEV can be traced to the assumption of a vanishing trilinear scalar coupling at the GUT scale. A small  $\mu_X$  can also be realized in a supersymmetric SO(10) context [30]. Regarding the second question, we recall the example of a non-supersymmetric SO(10) framework, which also contains the remnants of a larger  $E_6$  symmetry [27], where  $\mu_X$  is generated at two-loop level, but  $\mu_N$  is generated at a higher loop justifying its relative smallness. In our work, we do not advocate any specific GUT scenario to provide the dynamics that generates a small  $\mu_X$ . We treat  $\mu_X$  as an effective phenomenological parameter of unspecified origin, whose smallness derives its origin in some unknown hidden sector dynamics. We simply set its value to reproduce the correct light neutrino mass.

We now illustrate the pattern of neutrino masses with only one generation. In the  $\{\Psi_\nu, \Psi_N, \Psi_X\}$  basis, the  $(3 \times 3)$  neutrino mass matrix is given by

$$\mathcal{M} = \begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & M_N \\ 0 & M_N & \mu_X \end{pmatrix}, \quad (6)$$

yielding the mass eigenvalues ( $m_1 \ll m_{2,3}$ ):

$$m_1 = \frac{m_D^2 \mu_X}{m_D^2 + M_N^2}, \quad m_{2,3} = \mp \sqrt{M_N^2 + m_D^2} + \frac{M_N^2 \mu_X}{2(m_D^2 + M_N^2)}. \quad (7)$$

The important thing to observe here is that the lightness of the smallest eigenvalue  $m_1$  is due to the smallness of  $\mu_X$ . The other two eigenvalues ( $m_2$  or  $m_3$ ) can have a mass around 10 GeV, and their presence significantly influences the decay pattern of  $A_1$ .

We now compute the branching ratios of  $A_1$  into the invisible modes comprising of the  $\Psi_\nu$ ,  $\Psi_N$  and the  $\Psi_X$  states. Rigorously speaking, one should first diagonalize the mass matrix of Eq. (6) to determine the physical neutrino states. However, for our purpose it suffices to estimate the branching fractions of  $A_1$  into the  $\Psi_\nu \Psi_N$  and  $\Psi_N \Psi_X$  interaction states. Recall from Eqs. (1) and (5) that the decay of  $A_1$  into  $\Psi_\nu \Psi_N$  will depend on how large the doublet component of  $A_1$  is, i.e. on how large  $\cos \theta_A$  is, whereas the decay into  $\Psi_N \Psi_X$  will depend on the amount of  $A_S$  inside  $A_1$ , i.e. on the magnitude of  $\sin \theta_A$ . Below, we present the branching ratios into invisible modes normalised to the visible ones (neglecting, for simplicity, the phase space effects).

$$\frac{\text{Br}(A_1 \rightarrow \Psi_\nu \Psi_N)}{\text{Br}(A_1 \rightarrow f \bar{f}) + \text{Br}(A_1 \rightarrow c \bar{c})} \simeq \frac{m_D^2}{m_f^2 \tan^4 \beta + m_c^2}, \quad (8)$$

$$\frac{\text{Br}(A_1 \rightarrow \Psi_N \Psi_X)}{\text{Br}(A_1 \rightarrow f \bar{f}) + \text{Br}(A_1 \rightarrow c \bar{c})} \simeq \tan^2 \theta_A \frac{M_N^2}{m_f^2 \tan^2 \beta + m_c^2 \cot^2 \beta} \frac{v^2}{v_S^2}, \quad (9)$$

where  $v = \sqrt{v_u^2 + v_d^2} \simeq 174$  GeV. Notice that the dominant visible decay modes of  $A_1$  are  $f\bar{f}$  ( $f = \mu, \tau, b$ ) and  $c\bar{c}$ . Of course, the  $c\bar{c}$  mode would be numerically relevant if  $m_{A_1} < 2m_b$  and  $\tan\beta$  is small. Note that the branching ratio into  $\Psi_N\Psi_X$  dominates over that into  $\Psi_\nu\Psi_N$  for two reasons - firstly, there is a  $\tan^2\theta_A$  prefactor for the former which can be rather large if  $A_1$  has a dominant singlet component; secondly, if the  $m_f^2$  term in the denominator of the branching ratio expressions is numerically relevant, then the  $\Psi_\nu\Psi_N$  channel suffers a suppression by an additional  $\tan^2\beta$  factor.

For a numerical illustration, we make two choices of  $\tan\beta = (3, 20)$ , and fix  $\cos\theta_A = 0.1$ , which yield  $X_d = \cos\theta_A \tan\beta = (0.3, 2)$ . We recall that the upper limit on  $X_d$  for  $m_{A_1} < 8$  GeV in the minimal NMSSM has been obtained primarily from radiative  $\Upsilon$ -decays, and the limit is between 0.7 to 3.0 for  $\tan\beta = 50$ , while it is 30 or above for  $\tan\beta = 1.5$  [31]. A value of  $X_d = 2$  is in fact slightly above the upper limit for  $m_{A_1}$  in the range of 4 to 8 GeV. In the present scenario,  $A_1$  has a significant branching ratio into invisible modes which, in turn, considerably relax the upper bound on  $X_d$ . Here we do not choose a very large value of  $\tan\beta$  as that would increase the branching ratio of  $A_1$  into visible modes. The value of  $m_{A_1}$  is chosen to be somewhat larger than  $M_N$ , so that the phase space suppression, given by the factor  $\left(\left\{1 - \left(\frac{2m_f}{m_{A_1}}\right)^2\right\} / \left\{1 - \left(\frac{2M_N}{m_{A_1}}\right)^2\right\}\right)^{1/2}$ , is not numerically significant. We consider two values for  $M_N = (5, 30)$  GeV. The rationale behind choosing  $M_N = 5$  GeV is that it allows us to explore  $m_{A_1} < 10$  GeV, a regime where constraints from  $\Upsilon$ - and  $B$ -decays are particularly restrictive – see Table 1. On the other hand, the choice  $M_N = 30$  GeV implies that  $A_1$  is moderately heavy ( $m_{A_1} > 30$  GeV) which corresponds to the range where LEP and  $B$ -decay constraints are relevant. We display our results in Table 2. For numerical illustration, we have assumed  $v_S \sim \mathcal{O}(v)$ . The main conclusion is that if  $\cos\theta_A$  is small,  $A_1$  has a dominant singlet component (which is generally the case when  $A_1$  is light [4]), then for a reasonable part of the parameter space  $A_1$  can have a sizable invisible branching ratio which would weaken many of the constraints discussed in the beginning. However, it is important to stress that  $\cos\theta_A$  should not be excessively small, since in that case the purely singlet  $A_1$  would be completely decoupled from the visible sector.

	$\tan\beta = 20, \cos\theta_A = 0.1$		$\tan\beta = 3, \cos\theta_A = 0.1$	
$M_N$ (GeV)	5	30	5	30
$\text{Br}(A_1 \rightarrow \Psi_\nu\Psi_N)$	$7 \times 10^{-5}$	$3 \times 10^{-6}$	$4 \times 10^{-3}$	$1 \times 10^{-4}$
$\text{Br}(A_1 \rightarrow \Psi_N\Psi_X)$	0.7	0.9	$\sim 1$	$\sim 1$

Table 2: *Invisible branching ratios of the lightest NMSSM pseudoscalar for  $m_D = 10$  GeV,  $M_N = (5, 30)$  GeV, and  $\mu_X = 1$  eV.*

What about the CP-even Higgs mass limit? Since the  $ZZh$  coupling is diluted with respect to its SM value as a result of singlet admixture, the direct search (lower) limit on  $m_h$  from its non-observation will be lower than the SM limit of 114 GeV. A study based on the OPAL data from the LEP-2 run shows how the lower limit on  $m_h$  decreases (assuming the Higgs production via Higgs-strahlung process) as  $\xi \equiv \sigma(Zh)\text{Br}(h \rightarrow \text{invisible})/\sigma^{\text{SM}}(Zh)$  gets smaller than unity [32]. In our case, we have not only a mixing between the MSSM part of the CP-even Higgs and the singlet CP-even component, but also a sizable branching ratio of  $h \rightarrow A_1 A_1 \rightarrow \text{invisible}$ . The lower limit on the lightest CP-even Higgs mass will then decrease accordingly.

To conclude, we explore the possibility of having a very light (of order 10 GeV) pseudoscalar in the NMSSM, which has interesting consequences for CP-even neutral Higgs search at colliders, as well



as for facilitating dark matter annihilation. In this context, we have extended the minimal NMSSM with two additional gauge singlets carrying opposite lepton numbers for two specific reasons. On one hand, they provide a substantial invisible decay channel to the lightest pseudoscalar which helps relaxing or even evading some of the tight constraints from  $\Upsilon$ - and  $B$ -decays. On the other hand, they naturally set up the stage for implementing the inverse seesaw mechanism in order to generate light neutrino masses. What is phenomenologically interesting is that this can be done using order one neutrino Yukawa couplings and employing neutrino Dirac masses of a few tens of GeV. To account for the experimental values of the two mass squared differences and the three mixing angles of light neutrinos, one would of course have to extend the number of  $\hat{N}$  and  $\hat{X}$  superfields.

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## References

- [1] H. P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B **120** (1983) 346; J. M. Frere, D. R. T. Jones and S. Raby, Nucl. Phys. B **222** (1983) 11; J. P. Derendinger and C. A. Savoy, Nucl. Phys. B **237** (1984) 307; J. R. Ellis, J. F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D **39** (1989) 844; M. Drees, Int. J. Mod. Phys. A **4** (1989) 3635; U. Ellwanger, M. Rausch de Traubenberg and C. A. Savoy, Phys. Lett. B **315** (1993) 331; S. F. King and P. L. White, Phys. Rev. D **52** (1995) 4183; F. Franke and H. Fraas, Int. J. Mod. Phys. A **12** (1997) 479; M. Bastero-Gil, C. Hugonie, S. F. King, D. P. Roy and S. Vempati, Phys. Lett. B **489** (2000) 359.
- [2] C. Amsler *et al.* [Particle Data Group], Phys. Lett. B **667** (2008) 1, and <http://pdg.lbl.gov> for 2009 partial update.
- [3] M. Maniatis, Int. J. Mod. Phys. A **25** (2010) 3505.
- [4] U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. **496** (2010) 1.
- [5] R. Dermisek and J. F. Gunion, Phys. Rev. D **81** (2010) 075003.
- [6] S. Schael *et al.* [ALEPH Collaboration and DELPHI Collaboration and L3 Collaboration and OPAL Collaborations and LEP Working Group for Higgs Boson Searches], Eur. Phys. J. C **47** (2006) 547.
- [7] S. Schael *et al.* [ALEPH Collaboration], JHEP **1005** (2010) 049.
- [8] V. M. Abazov *et al.* [D0 Collaboration], Phys. Rev. Lett. **103** (2009) 061801.
- [9] R. Dermisek and J. F. Gunion, Phys. Rev. Lett. **95** (2005) 041801; Phys. Rev. D **73** (2006) 111701; Phys. Rev. D **75** (2007) 075019; Phys. Rev. D **76** (2007) 095006.
- [10] J. F. Gunion, D. Hooper and B. McElrath, Phys. Rev. D **73** (2006) 015011; D. Das and U. Ellwanger, JHEP **1009** (2010) 085; J. F. Gunion, A. V. Belikov and D. Hooper, arXiv:1009.2555 [hep-ph]; P. Draper, T. Liu, C. E. M. Wagner, L. T. M. Wang and H. Zhang, arXiv:1009.3963

- [hep-ph]; D. A. Vasquez, G. Belanger, C. Boehm, A. Pukhov and J. Silk, arXiv:1009.4380 [hep-ph].
- [11] F. Domingo and U. Ellwanger, JHEP **0712** (2007) 090; G. Hiller, Phys. Rev. D **70** (2004) 034018.
  - [12] W. Love *et al.* [CLEO Collaboration], Phys. Rev. Lett. **101** (2008) 151802.
  - [13] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **103** (2009) 081803.
  - [14] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **103** (2009) 181801.
  - [15] M. Drees and K. i. Hikasa, Phys. Rev. D **41** (1990) 1547.
  - [16] F. Domingo, U. Ellwanger and M. A. Sanchis-Lozano, Phys. Rev. Lett. **103** (2009) 111802.
  - [17] F. Domingo, arXiv:1010.4701 [hep-ph].
  - [18] G. Abbiendi *et al.* [OPAL Collaboration], Eur. Phys. J. C **23** (2002) 397.
  - [19] B. Aubert *et al.* [BaBar Collaboration], arXiv:0808.0017 [hep-ex]. See also an earlier analysis, R. Balest *et al.* [CLEO Collaboration], Phys. Rev. D **51** (1995) 2053.
  - [20] D. E. López-Fogliani and C. Muñoz, Phys. Rev. Lett **97** (2006) 041801; M. Chemtob and P.N. Pandita, Phys. Rev. D **73** (2006) 055012; A. Abada and G. Moreau, JHEP **0608** (2006) 044; A. Abada, G. Bhattacharyya, and G. Moreau, Phys. Lett. **B642** (2006) 503; N. Escudero, D.E. López-Fogliani, C. Muñoz, and R.R. de Austri, JHEP **0812** (2008) 099; J. Fidalgo, D.E. López-Fogliani, C. Muñoz, and R. Ruiz de Austri, JHEP **0908** (2009) 105; P. Ghosh and S. Roy, JHEP **0904** (2009) 069; A. Bartl, M. Hirsch, A. Vicente, S. Liebler and W. Porod, JHEP **0905**, 120 (2009); P. Ghosh, P. Dey, B. Mukhopadhyaya and S. Roy, JHEP **1005** (2010) 087.
  - [21] R. Kitano and K. y. Oda, Phys. Rev. D **61** (2000) 113001; M. Frank, K. Huitu, and T. Ruppel, Eur. Phys. J. **C52**, (2007) 413.
  - [22] D. G. Cerdeno, C. Munoz and O. Seto, Phys. Rev. D **79** (2009) 023510 [arXiv:0807.3029 [hep-ph]]; D. Das and S. Roy, Phys. Rev. D **82** (2010) 035002.
  - [23] R. N. Mohapatra and J. W. F. Valle, Phys. Rev. D **34** (1986) 1642.
  - [24] A. Abada, C. Biggio, F. Bonnet, M. B. Gavela and T. Hambye, JHEP **0712** (2007) 061.
  - [25] G. 't Hooft, Lecture given at Cargese Summer Inst., Cargese, France, Aug 26 - Sep 8, 1979.
  - [26] E. Dudas and C. A. Savoy, Acta Phys. Polon. B **33** (2002) 2547.
  - [27] E. Ma, Phys. Rev. D **80** (2009) 013013.
  - [28] F. Bazzocchi, arXiv:1011.6299 [hep-ph].
  - [29] F. Bazzocchi, D. G. Cerdeno, C. Munoz and J. W. F. Valle, Phys. Rev. D **81** (2010) 051701.
  - [30] P. S. B. Dev and R. N. Mohapatra, Phys. Rev. D **81** (2010) 013001.
  - [31] F. Domingo, U. Ellwanger, E. Fullana, C. Hugonie and M. A. Sanchis-Lozano, JHEP **0901** (2009) 061.
  - [32] K. Nagai, J. Phys. Conf. Ser. **110** (2008) 072028.